

Observation of a dynamical QCD string

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Mesons constructed from the quark propagators without the lowest-lying eigenmodes of the Dirac operator reveal not only restored $SU(2)_L \times SU(2)_R$ chiral and $U(1)_A$ symmetries, but actually a higher symmetry. All possible chiral and $U(1)_A$ multiplets for the states of the same spin are degenerate, i.e., the energy of the observed quantum levels does not depend on the spin orientation of quarks in the system and their parities. The quark-spin independence of the energy levels implies absence of the magnetic interactions in the system. The ultrarelativistic quark-antiquark system with only the color-electric interactions can be interpreted (or defined) as a dynamical QCD string.

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INTRODUCTION

With the static color charges one observes on the lattice a color-electric flux tube [1], which, if the distance between the static quarks is large, can be approximated as a string. This string is nondynamical (in the sense that its ends are fixed). In the light quark systems a crucial aspect is a relativistic motion of quarks at the ends of a possible string. There is no consistent theory of the dynamical QCD string with quarks at the ends. It is even apriori unclear whether such a picture has something to do with reality. In this case the chiral symmetry as well as its dynamical breaking should be relevant.

Both confinement and chiral symmetry breaking dynamics are important for the hadronic mass generation. Their interplay is responsible for a rather complicated structure of the hadron spectra in the light quark sector. A large degeneracy is seen in the highly excited mesons [2, 3], which is, however, absent in the observed spectrum below 1.8 GeV. The low-lying hadron spectra should be strongly affected by the chiral symmetry breaking dynamics.

In order to disentangle the confinement physics from the chiral symmetry breaking dynamics we remove on the lattice from the valence quark propagators the lowest-lying quasi-zero modes of the Dirac operator keeping at the same time the gluon gauge configurations intact [4–6]. Indeed, the quark condensate of the vacuum is related to a density of the lowest quasi-zero eigenmodes of the Dirac operator [7]:

$$\langle 0 | \bar{q}q | 0 \rangle = -\pi \rho(0). \quad (1)$$

We subtract from the valence quark propagators their lowest-lying chiral modes, which are a tiny part of the full amount of modes,

$$S_{RD(k)} = S_{Full} - \sum_{i=1}^k \frac{1}{\lambda_i} |\lambda_i\rangle \langle \lambda_i|, \quad (2)$$

where λ_i and $|\lambda_i\rangle$ are the eigenvalues and the corresponding eigenvectors of the Dirac operator. Given these reduced quark propagators we study the existence of hadrons and their masses upon reduction of the lowest-lying modes responsible for the chiral symmetry breaking.

LATTICE TECHNOLOGY

In contrast to refs. [4, 5], where a chirally not invariant Wilson lattice Dirac operator was employed, we adopt now [6] a manifestly chiral-invariant overlap Dirac operator [8]. The quark propagators were generously provided by the JLQCD collaboration [9–11]. For some technical details of the work we refer the reader to ref. [6].

OBSERVATIONS, SYMMETRIES AND THEIR STRING INTERPRETATION

At this first stage of the study we have investigated the ground state and the excited states of all possible $\bar{q}q$ isovector $J = 1$ mesons, i.e., $\rho(1^{--})$, $a_1(1^{++})$, $b_1(1^{+-})$. An exponential decay of the correlation function is interpreted as a physical state and its mass is extracted. The quality of the exponential decay improves with the reduction of the lowest modes. The evolution of masses of the ground and excited states upon reduction of the low-lying modes is shown in Fig. 1.

At the truncation energy about 50 MeV an onset of a degeneracy of the states ρ, ρ', a_1, b_1 , as well as a degeneracy of their excited states is seen. This degeneracy indicates a symmetry.

All possible multiplets of the $SU(2)_L \times SU(2)_R$ group for the $J = 1$ mesons are shown in Table 1 [2, 3]. In the chirally symmetric world there must be two independent ρ -mesons that belong to different chiral representations. The first one is a member of the $(0, 1) + (1, 0)$ representation. It can be created from the vacuum only by the operators that have the same chiral structure, e.g. by the vector-isovector current $\bar{q}\gamma^i \vec{\tau} q$. Its chiral partner is the

TABLE I: The complete set of $q\bar{q}$ $J = 1$ states classified according to $SU(2)_L \times SU(2)_R$. The symbol \leftrightarrow indicates the states belonging to the same representation R that must be degenerate in the chirally symmetric world.

R	mesons
$(0, 0)$	$\omega(I = 0, 1^{--}) \leftrightarrow f_1(I = 0, 1^{++})$
$(1/2, 1/2)_a$	$\omega(I = 0, 1^{--}) \leftrightarrow b_1(I = 1, 1^{+-})$
$(1/2, 1/2)_b$	$h_1(I = 0, 1^{+-}) \leftrightarrow \rho(I = 1, 1^{--})$
$(0, 1) \oplus (1, 0)$	$a_1(I = 1, 1^{++}) \leftrightarrow \rho(I = 1, 1^{--})$

axial vector meson, a_1 , that is created by the axial-vector isovector current. When chiral symmetry is restored this ρ -meson must be degenerate with the a_1 state. Another ρ -meson, along with its chiral partner h_1 , is a member of the $(1/2, 1/2)_b$ representation. It can be created only by the operators that have the same chiral structure, e.g. by $\bar{q}\sigma^{0i}\vec{\tau}q$.

In the real world (i.e. with broken chiral symmetry) each ρ -state (i.e. ρ and ρ') is a mixture of these two representations and they are well split. Upon subtraction of 10 lowest modes we observe two independent degenerate ρ -mesons. One of them couples only to the vector current and does not couple to the $\bar{q}\sigma^{0i}\vec{\tau}q$ operator. The other ρ meson - the other way around. This means that one of these degenerate ρ -states belongs to the $(0, 1) + (1, 0)$ multiplet and the other one is a member of the $(1/2, 1/2)_b$ representation.

A degeneracy of the $(0, 1) \oplus (1, 0)$ ρ -meson with the a_1 meson is a clear signal of the chiral $SU(2)_L \times SU(2)_R$ restoration. Consequently, a similar degeneracy should be observed in all other chiral pairs from Table 1.

The $U(1)_A$ symmetry transforms the b_1 state into the $(1/2, 1/2)_b$ ρ -meson [2, 3]. Their degeneracy indicates a restoration of the $U(1)_A$ symmetry. We conclude that simultaneously both $SU(2)_L \times SU(2)_R$ and $U(1)_A$ symmetries get restored.

The restored $SU(2)_L \times SU(2)_R \times U(1)_A$ symmetry requires a degeneracy of four mesons that belong to $(1/2, 1/2)_a$ and $(1/2, 1/2)_b$ chiral multiplets [2, 3], see Table I. This symmetry does not require, however, a degeneracy of these four states with other mesons, in particular with a_1 and its chiral partner ρ . We clearly see the latter degeneracy. This implies that there is some higher symmetry, that includes $SU(2)_L \times SU(2)_R \times U(1)_A$ as a subgroup. This higher symmetry requires a degeneracy of all eight mesons from the Table I.

All these eight mesons can be combined to a reducible $\bar{q}q$ representation, which is a product of two fundamental chiral quark representations [12]:

$$[(0, 1/2) + (1/2, 0)] \times [(0, 1/2) + (1/2, 0)] = \\ (0, 0) + (1/2, 1/2)_a + (1/2, 1/2)_b + [(0, 1) + (1, 0)]. \quad (3)$$

They exhaust all possible chiralities of quarks and antiquarks, i.e., their spin orientations, as well as possible spatial and charge parities for non-exotic mesons.

The observed degeneracy of all these eight mesons suggests that the higher symmetry, mentioned above, should combine the $U(1)_A$ and the $SU(2)_L \times SU(2)_R$ rotations in the isospin space with the $SU(2)_S$ spin-symmetry. The latter one is due to independence of the energy on the orientations of the quark spins. The higher symmetry group that combines all eight mesons from the Table 1 into one multiplet of dimension 16 should be $SU(2 \cdot N_f)$.

The quark-spin independence of the energy levels implies that there are no magnetic interactions in the system, i.e., the spin-orbit force, the color-magnetic (hyperfine) and tensor interactions are absent [13]. The energy of the system is entirely due to interactions of the color charges via the color-electric field and due to a relativistic motion of the system. We interpret (or define) such a system as a dynamical QCD string. Note a significant qualitative difference with the case of a motion of an electrically charged fermion in a static electric field or with a relative motion of two fermions with the electric charge. In the latter cases there exist spin-orbit and spin-spin forces that are a manifestation of the magnetic interaction in the system. In our case such a magnetic interaction is absent, which implies that both quarks are at rest with respect to the electric field (that moves together with quarks). It is this circumstance that suggests to interpret (define) our system as a dynamical QCD string.

The observed radial levels at the truncation energy 65 MeV, at which we see the onset of the symmetry, are approximately equidistant and can be described through the simple relation

$$E_{n_r} = (n_r + 1)\hbar\omega, \quad n_r = 0, 1, \dots \quad (4)$$

The extracted value of the radial string excitation quantum amounts to $\hbar\omega = (900 \pm 70)$ MeV. At the moment we cannot exclude, however, the quadratic relation $E_{n_r}^2 \sim (n_r + 1)$, because the excited level can be shifted up due to rather small finite lattice volume.

There is an interesting aspect of the dynamical QCD string that crucially distinguishes it from the Nambu-Goto open string. The energy of the the Nambu-Goto open bosonic string is determined by its orbital angular momentum L , $M^2 \sim L$. For the dynamical QCD string that contains chiral quarks at the ends the orbital angular momentum L of the relative motion is not conserved [14]. For instance, two orthogonal ρ -mesons at the same energy level are represented by the mutually orthogonal fixed superpositions of the S - and D -waves.

$$|(0, 1) + (1, 0); 1 \ 1^{--}\rangle = \sqrt{\frac{2}{3}}|1; {}^3S_1\rangle + \sqrt{\frac{1}{3}}|1; {}^3D_1\rangle, \\ |(1/2, 1/2)_b; 1 \ 1^{--}\rangle = \sqrt{\frac{1}{3}}|1; {}^3S_1\rangle - \sqrt{\frac{2}{3}}|1; {}^3D_1\rangle.$$

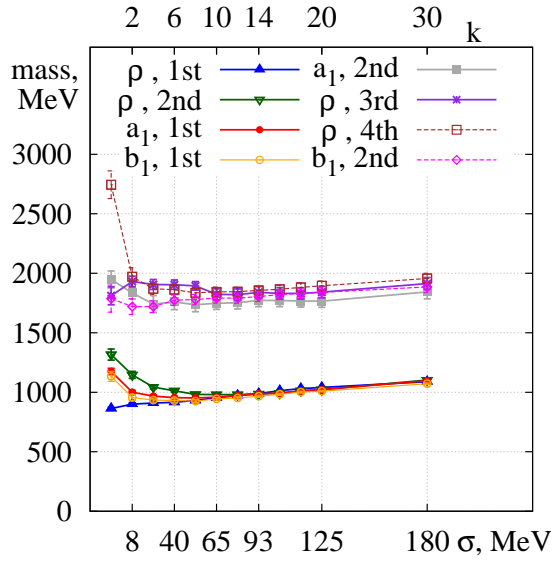


FIG. 1: Evolution of hadron masses under the low-mode truncation. Both the number k of the removed lowest eigenmodes as well as the corresponding energy gap σ are given.

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- [1] G. S. Bali, Phys. Rept. **343**, 1 (2001).
 - [2] L. Y. Glozman, Phys. Lett. B **587**, 69 (2004).
 - [3] L. Y. Glozman, Phys. Rept. **444**, 1 (2007).
 - [4] C. B. Lang and M. Schröck, Phys. Rev. D **84**, 087704 (2011).
 - [5] L. Y. Glozman, C. B. Lang and M. Schröck, Phys. Rev. D **86**, 014507 (2012).
 - [6] M. Denissenya, L. Y. Glozman and C. B. Lang, Phys. Rev. D **89**, 077502 (2014).
 - [7] A. Casher, Phys. Lett. B **83**, 395 (1979).
 - [8] H. Neuberger, Phys. Lett. B **417**, 141 (1998).
 - [9] S. Aoki *et al.* [JLQCD Collaboration], Phys. Rev. D **78**, 014508 (2008).
 - [10] S. Aoki *et al.* PTEP **2012**, 01A106 (2012).
 - [11] J. Noaki *et al.* [JLQCD and TWQCD Collaborations], Phys. Rev. Lett. **101**, 202004 (2008).
 - [12] T. D. Cohen and X. -D. Ji, Phys. Rev. D **55**, 6870 (1997).
 - [13] L. Y. Glozman, Phys. Lett. B **541**, 115 (2002).
 - [14] L. Y. Glozman and A. V. Nefediev, Phys. Rev. D **76**, 096004 (2007); Phys. Rev. D **80**, 057901 (2009).